

# Hot-Carrier Model for an Anomalous Exponent of Photoconduction

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Experiments often show that the photoconductance  $\sigma$  of a semiconductor system and the light intensity  $I$  are related by  $\sigma \sim I^\gamma$ . Conventional theories give a satisfactory explanation for  $\gamma = 1$  or  $\gamma = \frac{1}{2}$ , but anomalous exponents close to  $\gamma = \frac{3}{4}$  are often observed. This paper argues that there is a universal anomalous regime for which  $\gamma = \frac{3}{4}$  (or  $\gamma = \frac{2}{3}$  in two-dimensional samples), resulting from the kinetics of electron-hole recombination being controlled by Coulombic attraction. Because the local electric fields are extremely high, the theory appeals to the ‘hot-carrier’ model for transport in semiconductors.

**Keywords:** Photoconduction, recombination

## 1. Introduction

Many semiconductor devices rely on excess charge carriers which are either injected or else created by absorbing photons (Sze, 1981). Electron-hole recombination is therefore a process which is central to understanding the dynamics of these systems. The radiative and non-radiative mechanisms for recombination have been extensively studied, starting with seminal work of Shockley and Read (1952). In this paper I shall argue that there is a universal regime for the kinetics of recombination which has not been described in earlier work. The evidence comes from considering an anomalous exponent describing the response of photoconductors, where the photoconductance  $\sigma$  and the light intensity  $I$  are typically related by  $\sigma \sim I^\gamma$ : values close to  $\gamma = \frac{3}{4}$  are frequently observed, whereas conventional theory only gives a satisfactory account for  $\gamma = 1$  or  $\gamma = \frac{1}{2}$ . I argue that accounting for the role of Coulomb interaction in the kinetics of electron-hole recombination leads to an explanation for the anomalous exponent,  $\gamma$ . A version of this model which uses the ‘hot-carrier’ concept (discussed in Ziman, 1972), as opposed to linear response, provides the most satisfying description and the best fit to experimental data. This model gives  $\gamma = \frac{3}{4}$  in three dimensions, and  $\gamma = \frac{2}{3}$  for ‘two-dimensional’ systems, where the photoresponsive layer is very shallow.

The photocurrent excited in a semiconductor might be intuitively expected to be proportional to the light intensity, and there are reasonable models (discussed below) which do imply a linear dependence. A simple and plausible argument however leads to a different conclusion. Consider a simple photocell consisting of a semiconductor (undoped) with two electrodes attached. Light falling on the device creates electron-hole pairs (at a rate which is proportional to the intensity of radiation,  $I$ ). These mobile carriers allow the device to conduct electricity, with a conductivity

$\sigma$ , which is proportional to the density of charge carriers. If the density of mobile electrons is  $n_e$  and the density of holes is  $n_h$ , the conductivity is

$$\sigma = e(D_e n_e + D_h n_h) \quad (1.1)$$

where  $D_e$  and  $D_h$  are the electron and hole diffusion coefficients, and  $e$  is the electron charge. Because the semiconductor is assumed to be intrinsic, the number of electrons equals the number of holes, so that  $n_e = n_h \equiv n$ . The rate of change of the excitation density  $n$  contains a source term, proportional to the intensity, and a sink term, which represents the rate for electron-hole recombination:

$$\frac{dn}{dt} = AI - f(n) \quad (1.2)$$

where  $A$  is a constant and  $f(n)$  is the rate for recombination of electron-hole pairs at density  $n$ . It is natural to invoke the law of mass action, and to propose that the rate of recombination is proportional to the product of the number density of electrons and of holes: that is the recombination term is  $dn_e/dt = dn_h/dt = -Cn_en_h$  (for some constant  $C$ ), so that  $f(n) = Cn^2$ . In this case the equilibrium concentration  $n$  and consequently the conductance are proportional to  $\sqrt{I}$ .

Thus there appear to be two plausible models for the recombination rate, which are distinguished by looking at the exponent  $\gamma$  for a power-law dependence of conductance  $G$  upon light intensity:

$$G \sim I^\gamma. \quad (1.3)$$

The model described above indicates  $\gamma = \frac{1}{2}$ . The intuitive expectation that  $\gamma = 1$  is consistent with a picture where the electrons and holes become immobilised on ‘recombination centres’, which exist independently of the illumination.

It is not immediately clear which value of  $\gamma$  is correct, and it is necessary to turn to experiment to decide the issue. Both values,  $\gamma = 1$  and  $\gamma = \frac{1}{2}$  are observed in different experiments, corresponding respectively to mobile charge carriers being destroyed by ‘monomolecular’ or ‘bimolecular’ processes. But a surprising aspect of the experimental literature is that intermediate values of  $\gamma$  (hereafter termed *anomalous*) are often observed over several decades of intensity.

The existence of anomalous exponents has been discussed by Rose (1963), but his model is unsatisfying because it assumes that there is a set of impurity states in the band gap with a density of states which is an exponential function of energy. Careful studies of Lifshitz tails pioneered by Halperin and Lax (1966) indicate that the density of states has a more complex functional dependence, such as an exponential of a power of energy (measured from the band edge): see Thirumalai (1986). Because of these results it seems implausible to use an assumption that the density of states is a simple exponential in order to explain an anomalous response which has a well-defined exponent over decades of intensity. The model described by Rose is also vulnerable to a more fundamental objection, namely that it relies on an artificial demarcation between electron states which are traps and those which are recombination centres. The theory proposed by Rose indicates that  $\gamma$  should have a monotonic variation with temperature. This hypothesis is not compatible with experimental observations by Wronski and Daniel (1981) where the temperature was varied over a wide range: the exponent of the anomalous regime was found to

be insensitive to temperature (which is consistent with the ‘universality’ hypothesis developed in this paper).

Reading the literature suggests that where anomalous exponents are reported, they are frequently close to  $\frac{3}{4}$ , particularly in cases where a convincing fit to a power-law is demonstrated over two or more decades of intensity. The works which I have become aware of are discussed in section 2. Allowing for errors in fitting the exponent, the experimental results suggest that  $\gamma = \frac{3}{4}$  could be a ‘universal’ exponent for the anomalous regime.

The samples used in experiments where the anomalous response is observed often require specialist equipment for their preparation, and the reported values of  $\gamma$  are not always consistent. In view of the very surprising nature of the anomalous photoconductive response, it is desirable to characterise the effect in a readily accessible system. As well as reviewing the experimental literature, section 2 will discuss the intensity dependence of the conductance of some commercially available photocells. An anomalous photoconductive response with  $\gamma = \frac{3}{4}$  was observed in three of the four samples. The remaining one showed an anomalous regime with  $\gamma = \frac{2}{3}$ .

Section 3 discusses the theory for the rate of recombination. It is pointed out that the law of ‘mass action’ assumes implicitly that electrons and holes find each other ( $\gamma = \frac{1}{2}$  case) or a recombination trap ( $\gamma = 1$  case) by diffusion. Because the electrons and holes are electrically charged, there is a regime in which their attraction ensures that they collide more rapidly than if their motion was undirected diffusion. I show how this leads to a model for the anomalous exponent  $\gamma$ . The value of  $\gamma$  predicted by this model depends upon how the carrier drift velocity  $v$  varies with the electric field  $\mathcal{E}$ . The natural assumption is to apply linear response theory, and to assume that  $v \propto \mathcal{E}$ . However, the electric field due to the electrostatic attraction of the carriers is extremely large. It is known that for very large electric fields the motion of charge carriers becomes very insensitive to the electric field (Ziman, 1972). This is accounted for by assuming  $v \propto \mathcal{E}^\beta$ , with  $\beta$  very small. In the limit as  $\beta \rightarrow 0$ , the expression for the anomalous exponent gives  $\gamma = \frac{3}{4}$  in three dimensions and  $\gamma = \frac{2}{3}$  in two dimensions.

## 2. Experimental evidence

### (a) Discussion of earlier investigations

Anomalous photoconductive response has been observed in a wide variety of semiconductors and semiconductor based devices. The results which I have found reported are summarised in table 1. Some authors (Bakr, 2002, Kaplan and Kaplan, 1998, 2002, Spear *et al*, 1974) report a broad range of values of  $\gamma$ , with the exponent depending upon temperature, or other parameters. In the remaining cases where either a single value of  $\gamma$  or a narrow range is reported, its value appears to lie in the range 0.7 – 0.8, apart from three ‘outliers’ which are discussed below. Small deviations of the exponent fitted in experiments can result from various effects, such as the limits of the intensity range extending into regions where the anomalous regime is breaking down, or the dielectric properties of the materials being modified by the concentration of mobile carriers. For this reason, it is unlikely that differences in these values of  $\gamma$  are significant. The experimental literature is consistent with

Table 1. *Summary of previously published anomalous exponents,  $\gamma$ .*

Source	Material	$\gamma$	Remarks
Arene & Baixeras (1984)	$a$ -Si H	0.78	4 decades
Arene & Baixeras (1984)	$a$ -Si H	0.55	3.5 decades
Bakr (2002)	As <sub>2</sub> Sb <sub>3</sub> film	0.58 – 0.71	1.5 decades
Jie <i>et al</i> (2006)	Cd S nanowires	0.74-0.77	1.5 decades
Kaplan & Kaplan (1998)	As <sub>2</sub> Se <sub>3</sub>	0.44 – 0.86	2 decades, freq. varied
Kaplan & Kaplan (2002)	$a$ -Se	0.61-0.96	2 decades, freq. varied
Kind <i>et al</i> (2002)	Zn O nanowires	0.8	4 decades
quoted in Rose (1963)	Sb <sub>2</sub> S <sub>3</sub> (Vidicon cameras)	0.68	‘several decades’
Spear <i>et al</i> (1974)	$a$ -Si	0.54-0.9	method unclear
Wronski & Daniel (1981)	$a$ -Si H	0.83	2-3 decades, T varied
Wronski & Daniel (1981)	$a$ -Si H	0.7	3 decades

the view that there could be a robust regime with an exponent close to the middle of the range 0.7 – 0.8, while different exponents may be observable in other systems.

There are three ‘outliers’, where a single value of  $\gamma$  is quoted which is not in the range 0.7 – 0.8. Wronski and Daniel show results which are consistent with  $\gamma = 0.83$  over a range of temperatures. Rose (1963) reports that Vidicon camera tubes (which used Sb<sub>2</sub>S<sub>3</sub>) give  $\gamma = 0.68$  over ‘several decades’. This value is sufficiently far from the other exponents that it might have a different origin. Also Arene and Baixeras (1984) show a plot with  $\gamma = 0.55$  over nearly four decades. In this case, the exponent is so close to the bimolecular case  $\gamma = \frac{1}{2}$  that it is questionable whether this is an anomalous response.

The theory presented by Rose (1963) predicts that  $\gamma$  has the following dependence upon the temperature  $T$ :

$$\gamma = \begin{cases} \frac{T_0}{T+T_0} & T \leq T_0 \\ \frac{1}{2} & T > T_0 \end{cases} \quad (2.1)$$

where  $T_0$  is an effective temperature characterising the exponential density of states in the band gap. A more recent paper developing this approach is Schellenberg and Kao (1988). Wronski and Daniel (1981) examined the intensity dependence of the photocurrent in hydrogenated amorphous silicon at a variety of temperatures ranging from 147K to 267K. Figure 5 of their paper demonstrates that the exponent  $\gamma$  is close to 0.83 and independent of  $T$  at lower intensities, crossing over to 0.5 at higher intensity. Their data are, therefore, incompatible with (2.1).

There are some published results which contradict the evidence in Wronski and Daniel (1981), but these works do not present a coherent picture. Papers by Spear *et al* (1974), Kaplan and Kaplan (1998) and Bakr (2002) report that  $\gamma$  does depend upon temperature, but their exponents vary erratically and do not show agreement with (2.1). It is not clear how the intensity dependence of the exponents obtained by Spear *et al* were extracted from the experimental data.

In table 1 I have not included reports (such as Bube, 1957, Bube and Lind, 1958) where anomalous photoconductive response was observed but where the samples displayed non-Ohmic behaviour.

*(b) Investigation on commercial photocells*

In discussing the earlier experimental work I have pointed out that much of the published experimental data are consistent with a previously un-remarked universality, where  $\gamma \approx \frac{3}{4}$ . Some of the other literature suggests, however, that the values of the anomalous exponent are highly variable and may be erratic and hard to reproduce. By definition, universal behaviour does not shy away from observation. It might be observable in a simple experiment with a generic semiconductor sample. With this motivation, I measured the response of four different commercially available light-dependent resistors at room temperature (14 – 17°C). The photocells, samples 1-4, were obtained from Maplin Electronics. Their catalogue numbers were: sample 1 - N56AY, sample 2 - N57AY, sample 3 - N53AY, sample 4 - N46AY. The specifications do not include the composition of the sensitive semiconductor layer, but that is not relevant to the question of whether the anomalous exponent is universal.

The intensity of light from a Tungsten bulb was reduced by stopping down the aperture of a light box, and by increasing the distance between the aperture and the photocell. Some of the higher intensity data used an unshielded 15 W bulb, varying the distance from the sample. The response of each photocell was found to be Ohmic for both low and high light intensities, and its dark current was negligible. All of the data plotted in figure 1 used a potential of 3 V.

The intensity was varied over more than four decades. Figure 1 shows logarithmic plots of conductance  $G$  (in units of  $10^{-6}\Omega^{-1}$ ) versus intensity  $I$ . The light intensity scale is in arbitrary units, chosen so that the 15 W pearlescent bulb at 1 m gave an intensity  $I = 4 \times 10^{-3}$ . For each sample there was clear evidence for an anomalous regime. Sample 1 shows a monomolecular regime ( $\gamma = 1$ ) at low intensity, crossing over to an anomalous regime at higher intensity for which  $\gamma = \frac{3}{4}$  gives an excellent fit over two decades. Sample 2 showed an anomalous regime with  $\gamma = \frac{3}{4}$  spanning nearly two decades at intermediate intensities between regimes with  $\gamma = 1$  and  $\gamma = \frac{1}{2}$ . Sample 3 showed an anomalous regime for which  $\gamma = \frac{2}{3}$  provides an excellent fit over three decades, crossing over to bimolecular behaviour ( $\gamma = \frac{1}{2}$ ) at high intensities. Sample 4 shows an anomalous regime for which  $\gamma = \frac{3}{4}$  provides a good fit, crossing over to  $\gamma = \frac{1}{2}$  at high intensity.

Three out of four samples showed anomalous behaviour with  $\gamma = \frac{3}{4}$ , providing strong support for the hypothesis that this is a ‘universal’ phenomenon. It is noteworthy that the other sample produced an exponent very close to that reported by Rose (1963) for the Vidicon camera. All four samples showed crossover to other regimes, similar to that seen in figure 5 of Wronski and Daniel (1981).

### 3. Theory

*(a) A critique of the law of mass action and an alternative theory*

My own experiments and much of the previously published work is consistent with the view that there are three different regimes which can be observed in different systems, or over different ranges of intensity within the same system. In the order where they might be expected to be observed as the light intensity is increased, these are:

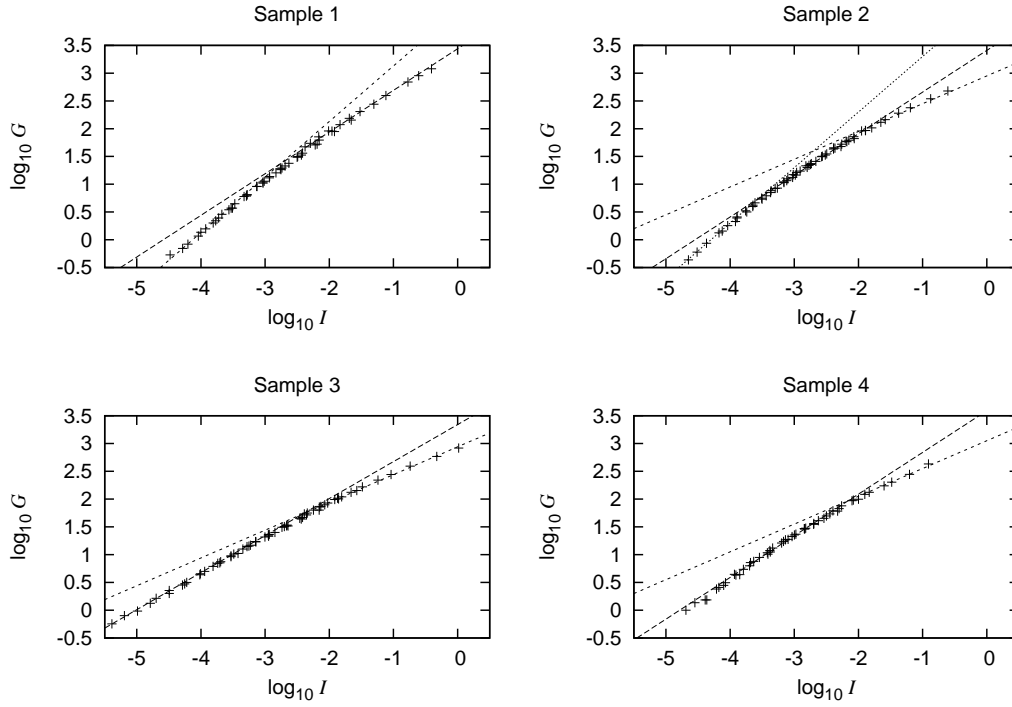


Figure 1. Plot of  $\log_{10} G$  (the conductance  $G$  measured in units of  $10^{-6} \Omega^{-1}$ ) versus  $\log_{10} I$  (with the intensity  $I$  in arbitrary units). Each sample showed an anomalous regime. The slopes of the linear fits are as follows. Sample 1:  $\gamma = 1, \frac{3}{4}$ , Sample 2:  $\gamma = 1, \frac{3}{4}, \frac{1}{2}$ , Sample 3:  $\gamma = \frac{2}{3}, \frac{1}{2}$ , Sample 4:  $\gamma = \frac{3}{4}, \frac{1}{2}$ .

1. *Monomolecular regime*. Here  $\gamma = 1$ , which is consistent with carriers being trapped on a fixed number of recombination centres. In the low-intensity limit it is expected that a carrier will always encounter such a trap before encountering a carrier with the opposite sign.
2. *Anomalous regime*. The origin of this will be considered below. In cases where a crossover has been observed as intensity is varied, this regime appears to occur in between the monomolecular and bimolecular regimes.
3. *Bimolecular regime*. Here  $\gamma = \frac{1}{2}$ , which is consistent with carriers being removed by electron-hole recombination at a rate proportional to the product of the electron and hole concentrations.

As the terms *monomolecular* and *bimolecular* suggest, the exponents  $\gamma = 1$  and  $\gamma = \frac{1}{2}$  are both explained by invoking the law of mass action, as outlined in the introduction. The law of mass action is applicable in situations where the reactive species move randomly by diffusion, without any interaction occurring until they are in contact. This idealisation is realistic for chemical reactions in which there is no electrostatic interaction. It may also be valid when the species are charged but when the rate constant for the chemical reaction is small, so the species reach a local quasi-equilibrium.

In the case of electron-hole recombination, there is Coulombic attraction between electrons and holes. This can pull them together, so that the time required for a collisions to occur may be less than it would be for randomly diffusing particles. The effect of recombination events on the density  $n$  may be written  $dn/dt = n/\tau$ , where  $n$  is the number of either species, and  $\tau$  is the lifetime for an electron or hole to survive without collisions. The following discussion uses scaling arguments rather than detailed calculations of the electron dynamics, but these arguments are sufficient to make a precise determination of the anomalous exponent.

If the collisions occur due to random diffusion of electrons and holes, then the lifetime  $\tau$  will be inversely proportional to the density of the other species, so that  $dn/dt \propto n^2$ . If the Coulombic force is significant, the timescale for collision will be the time required for an electron and a hole to be pulled together by their mutual electrostatic attraction, and  $\tau$  may have a different dependence upon the carrier density,  $n$ . In this case  $\tau$  is expected to be proportional to the characteristic distance  $L$  that an electron and a hole have to travel in order to make a collision. The timescale is also inversely proportional to the characteristic velocity  $v$  for relative motion of the electrons and the holes:  $\tau \sim L/v$ . If the density of carriers is  $n$ , we clearly have  $L \sim n^{-1/d}$ , where  $d$  is the effective spatial dimension of the photosensitive region penetrated by the radiation. (That is, we set  $d = 3$ , unless the depth is small compared to  $L$ , when we would use  $d = 2$ ). The velocity  $v$  will depend upon the local electric field  $\mathcal{E}$  driving the drift towards collision. The charges  $\pm e$  of the electrons and holes are independent of the density, but the lengthscale  $L$  depends upon the carrier density so that  $\mathcal{E} \sim L^{-1} \sim n^{1/d}$ .

In order to determine how  $\tau$  depends upon  $n$ , it remains to specify how the typical drift velocity  $v$  depends upon the electric field  $\mathcal{E}$ . There are two possibilities:

1. At first sight it seems natural to invoke linear response theory, and to propose that  $v \propto \mathcal{E}$ . It should, however, be noted that because the lengthscale  $L$  is very small, the electric field  $\mathcal{E}$  will be very large. For example, the electric field due to an electron at a distance of  $10^{-8}$  m is of order  $10^7$  V m<sup>-1</sup>. These electric fields are so high that the applicability of linear response theory is questionable.
2. It is argued that ‘hot-carriers’ in semiconductors have a damping rate due to interaction with phonons which rises very steeply with energy (Ziman 1972). As a consequence, the velocity of mobile carriers in semiconductors becomes insensitive to the applied electric field, and may even approach a limiting value as the electric field increases (Ferry 1975, Arora 1984). These hot carrier effects are considered to be significant for electric fields in excess of  $10^5$  V m<sup>-1</sup>.

The physics underlying the hot carrier model will be discussed in greater detail below, after considering its consequences. In the following it will be assumed that the relation between carrier velocity and electric field may be described by a power-law, such that

$$v \sim \mathcal{E}^\beta . \quad (3.1)$$

The linear response model corresponds to setting  $\beta = 1$ . The hot carrier model, where the velocity is quite insensitive to the microscopic electric field, is consistent with either assigning a very small value to  $\beta$ , or else taking the limit as  $\beta \rightarrow 0$  in

the final result. Before discussing the value of  $\beta$ , let us consider the formula for the anomalous exponent. We have a recombination rate of the form

$$\frac{dn}{dt} \sim -\frac{nv}{L} \sim nn^{1/d}\mathcal{E}^\beta \sim n^{(1+\frac{1}{d}+\frac{\beta}{d})}. \quad (3.2)$$

Since the rate of recombination is proportional to the flux  $I$ , we find  $n \sim I^\gamma$  with

$$\gamma = \left(1 + \frac{1}{d} + \frac{\beta}{d}\right)^{-1} = \frac{d}{d+1+\beta} \quad (3.3)$$

so that this theory is consistent with the photoconductivity having a power-law dependence upon  $I$  with a non-trivial exponent. Equation (3.3) is the central theoretical result of this paper. In three dimensions the exponent takes the value  $\gamma = \frac{3}{5}$  for the first model considered above, where the drift velocity of the carriers is determined by linear response theory. In the second version of the model, where the hot-carrier theory is applied and the limit  $\beta \rightarrow 0$  is taken, the exponent is  $\gamma = \frac{3}{4}$ , which is consistent with most of the experimental results. For  $d = 2$ , the hot carrier model gives  $\gamma = \frac{2}{3}$ , which is consistent with the results for my own sample 3, with the exponent quoted by Rose (1963) for Vidicon cameras, and with the mid-range for the results of Bakr (2002) on  $\text{As}_2\text{Sb}_3$  films.

It should be noted that while the non-linear ‘hot-carrier’ effect influences the timescale for electron-hole combination, the response to the small externally applied field is determined by equation (1.1) and is Ohmic.

Other authors (see Schlagenotto, Meader and Gerlach, 1974, Hangleiter, 1993) have suggested that Coulomb interactions can influence the rate of electron-hole recombination, but these works consider electron-hole correlations which will modify the constant  $C$  in the mass-action formula for the collision rate,  $dn_e/dt = C n_e n_h$ . Their calculations are only relevant to the bi-molecular process.

#### (b) Discussion of the hot carrier theory

Now consider the response of an electron or hole to the local electric field. Note that because we consider the microscopic structure of the electric field on a scale of the separation of the carriers, the field is not screened by polarisation effects. The samples where the anomalous response is observed are usually amorphous or highly disordered microcrystalline systems. The following discussion will assume that the carriers and the phonons are moving in a highly disordered environment.

The energy available from the electrostatic interaction potential is much higher than the thermal energy. The electrons and holes can therefore become ‘hot’ carriers’, because their energies relative to the band edge may be much higher than the thermal energy scale.

Studies of spectral hole-burning in disordered semiconductors (Hegarty and Sturge, 1985) indicate that the motion of electrons and holes is strongly damped by the effects of phonons. Because of the energy loss to phonons, the electrons and holes are expected to remain close to the bottom of the conduction band or to the top of the valence band, respectively, despite having energies which are higher than the thermal energy.

The physics of hot-carrier transport is extremely complex, and in the following discussion issues such as valley degeneracy will be ignored. The objective is to give



a motivation for assigning a small value to  $\beta$  in (3.1), based on a reasonable but simplified model for the interaction of charge carriers and phonons.

The electric field  $\mathcal{E}$  will cause an electron (or hole) to accelerate to a velocity  $v$  corresponding to an energy  $E = \frac{1}{2}mv^2$ , where  $m$  is the effective mass of the carrier (we assume that the electrons or holes remain close to bottom or top of their bands, where the particle dispersion relations are approximately parabolic and the phonon dispersion can be neglected). A particle with velocity  $v$  will then gain energy from the electric field at a rate  $\dot{E} = e\mathcal{E}v \sim \mathcal{E}E^{1/2}$ . This acceleration is opposed by the tendency of the motion of the charged particle to excite phonons: a particle with energy  $E$  will dissipate energy to phonons at a rate which will be written  $\dot{E} = R(E)E$ , where  $R(E)$  is a rate coefficient. The energy of a carrier accelerated by a constant electric field is therefore given by the solution of the energy balance equation

$$\frac{dE}{dt} = e\mathcal{E}\sqrt{2E/m} - R(E)E. \quad (3.4)$$

The hot carrier picture asserts that rate coefficient  $R(E)$  increases very rapidly with energy  $E$ . Various arguments and calculations have been advanced to support this claim (Ziman 1972, Ferry 1975, Arora 1984). In the following I present a simplified but plausible model, which yields  $R(E) \sim E^{9/2}$ .

The loss of energy by an electron to the phonons can be seen as analogous to spontaneous emission from an atomic state, with the phonons playing a role analogous to the photons in electrodynamics. Recall the expression for the rate  $R_{nm}$  of spontaneous emission from an atomic level with energy  $E_n$  to a level with lower energy,  $E_m$ :

$$R_{nm} = \frac{4\alpha}{3c^2} |M_{nm}|^2 \omega_{nm}^3 \quad (3.5)$$

where  $\omega_{nm} = (E_n - E_m)/\hbar$  is the frequency of the photon which is emitted, and  $M_{nm}$  is the corresponding dipole matrix element ( $c$  is the speed of light and  $\alpha$  is the fine structure constant). This expression can be adapted to describe the spontaneous decay of electronic states by emission of phonons. For simplicity it will be assumed that the matrix elements  $M_{nm}$  for the electron-phonon coupling in the highly disordered environment are random numbers, with statistics which are independent of the energies  $E_n$  and  $E_m$ . The rate of emission into any given state of lower energy is therefore proportional to  $\Delta E^3$ , where  $\Delta E = E_n - E_m$ . The number of available electronic states with energy below  $E$  is  $N(E) \sim E^{3/2}$ . Upon integrating over the lower energy, the rate coefficient scales as  $R(E) \sim E^{9/2}$ . The steady-state solution of equation (3.4) is therefore  $\mathcal{E} \sim E^5$ , so that the terminal velocity is

$$v \sim \mathcal{E}^{1/10} \quad (3.6)$$

that is  $\beta = \frac{1}{10}$ . This supports the assertion of the hot-carrier model, that the electron or hole drift velocity becomes highly insensitive to the electric field. Some authors (for example, Ferry 1975, Arora 1984) have gone further, and have suggested that the drift velocity approaches an asymptote as the electric field increases. This corresponds to setting  $\beta = 0$  in (3.3).

## 4. Conclusions

The theoretical understanding of the anomalous photoconductive response has been based upon arguments due to Rose (1963), which make a questionable assumption about the density of trapping states in the band gap. This theory predicts that the anomalous exponent  $\gamma$  is temperature dependent, in contradiction with experiments by Wronski and Daniel (1981).

This paper has proposed an alternative explanation of the anomalous response, based upon the idea that the kinetics of electron-hole recombination are driven by Coulomb attraction. The recombination time is then determined by the distance carriers have to travel, and their velocity  $v$ , which is assumed to have a power-law dependence upon the local electric field  $\mathcal{E}$ :  $v \sim \mathcal{E}^\beta$ . This leads to a simple expression for the anomalous exponent, equation (3.3). Because the electric fields are extremely high it is not appropriate to use linear response theory. In the hot-carrier picture for transport at high electric fields, the drift velocity has very weak dependence upon the electric field, which justifies setting  $\beta = 0$  in equation (3.3).

This model leads to the prediction that the anomalous regime has universal exponents of  $\gamma = \frac{3}{4}$  (three dimensions) or  $\gamma = \frac{2}{3}$  (for thin samples which are effectively two-dimensional). In support of this universality hypothesis, I have pointed out that many of the reported values of  $\gamma$  for the anomalous regime are close to one or other of these numbers. Furthermore, experiments on four commercially available photocells showed anomalous regimes with  $\gamma = \frac{3}{4}$  in three cases, and  $\gamma = \frac{2}{3}$  in the remaining case.

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